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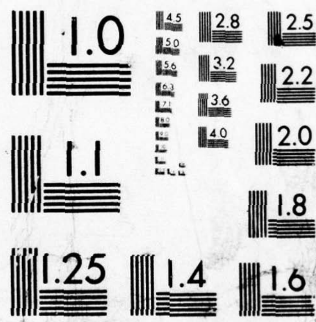

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Light Absorption by Ion Acoustic Turbulence in Laser-Produced Plasmas

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Plasma Physics Division

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A review of experiments on light absorption by laser-produced plasmas is presented and the need for another efficient absorption mechanism besides resonant absorption is clearly demonstrated. Laser light absorption by ion acoustic fluctuations is proposed as the missing absorption mechanism. Calculations of this absorption mechanism in a magnetized plasma are then presented.

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CONTENTS

LIGHT ABSORPTION BY ION ACOUSTIC TURBULENCE IN LASER-PRODUCED PLASMAS

I INTRODUCTION

The problem of how laser light is absorbed in a laser produced plasma is of crucial importance to the laser fusion program, as well as being a very interesting scientific problem in its own right. This paper studies anomalous absorption of laser light by ion acoustic turbulence in the underdense plasma. It is a continuation of our earlier work on this subject[1], which studied laser light absorption in an unmagnetized plasma by self generated ion turbulence.[2] In this paper we discuss this mechanism further, and also calculate absorption by ion acoustic turbulence in a *magnetized plasma*.[3]

A crucial characteristic of this absorption mechanism is that it occurs in the underdense plasma where gradient scale lengths are expected to be fairly long, rather than at the critical density where the gradient scale length is expected to be short. It is interesting to note that resonant absorption, which occurs at the critical surface, works best if the density gradient scale lengths are short. Therefore, since these two absorption mechanisms occur at different spatial positions, there is no reason why they cannot both occur simultaneously. A schematic of the expected plasma density profile, with the regions of resonant absorption, and absorption by ion acoustic turbulence is shown in Fig. 1. The units are free space wavelength of the laser light.

Another interesting aspect of absorption by ion acoustic turbulence is that a single fluid theory is valid for all electrons. Although the classical mean free path of an energetic electron is long compared to a gradient scale length, this mean free path is anomalously reduced in the

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presence of turbulence and/or a magnetic field.[1,2,3] Thus it is possible that a single fluid theory, with anomalous transport, but without special treatment for energetic electrons, can give an adequate description of electron thermal transport in a laser produced plasma.

Section II of this paper reviews recent data on laser light absorption at a number of different laboratories and shows that resonant absorption cannot account for more than about one fourth to one half of the total absorption. Section III discusses anomalous absorption in an unmagnetized plasma, while Sections IV and V discuss absorption in a magnetized plasma. Section VI shows how our results scale with laser wavelength, and Section VII discusses how to solve for the dynamics of the problem instead of the steady state.

II THE NEED FOR ANOTHER ABSORPTION MECHANISM

Before detailing our proposed absorption mechanism, we will briefly review some laser plasma experiments and then show that the three absorption mechanisms currently considered, inverse Bremsstrahlung, resonant absorption and parametric instabilities are not sufficient to explain the measured laser light absorption. There are now many different experiments, from many different laboratories which show fractional light absorption of order 50% or more.[4–11] Several of these measurements show that the fractional absorption is not strongly dependent on either tilt angle of the target[4,11] or polarization of the laser light.[4] Some of these,[7,11] and also other experiments[12,13] measure the density gradient scale length in the underdense region. For neodymium laser experiments, underdense scale lengths are typically between 10 and 60 microns,[7,12,13] where here the underdense scale length is defined as the distance between the critical and quarter critical surfaces. For CO_2 laser experiments, they are typically larger than 50 microns.[11]

Experiments at the Naval Research Laboratory have consistently shown high light absorption on experiments without prepulses. For instance Ref 4 shows absorption between fifty

and seventy percent which is insensitive to polarization and target tilt angle for tilt angles less than about 50°-60°. Furthermore, experiments in the Soviet Union[6,7] with a neodymium laser have also shown fifty to seventy percent absorption.

Experiments with CO_2 laser produced plasmas show high absorption too. For instance the experiments of Donaldson et al[9] show absorption of around 90%. The experiments of Richardson[11] show absorption of about 50 percent when the target is in the focal plane of the lens. Furthermore, this absorption is relatively independent of the target tilt angle for tilt angles of less than about thirty degrees. Finally, the experiments of Giovanielli et al[10] also show that the fractional absorption is roughly 45%.

One very interesting series of experiments is that reported by Van Kessel et al.[5] They showed that for a neodynium laser produced plasma, the absorption depends only weakly on laser light irradiance, but depends quite strongly on the focal conditions. If the target is in the focal plane of the lens, the absorption maximizes at a value of between forty five and sixty percent. As the target was moved in front of or in back of the focal plane, the fractional light absorption fell off. A similar result was found in a recent eperiment at the Naval Research Laboratory.[15] There, the laser irradiance was varied between 4×10^{14} W/cm^2 and 2×10^{16} W/cm^2 at focus. The absorption was relatively constant at between fifty and sixty percent. However it fell off when the target was moved out of the focal plane. For instance two shots at the same intensity, but different focal conditions had markedly different fractional absorption. Finally, we note that this same dependence of fractional absorption on focal conditions was found for a CO_2 laser produced plasma also.[11]

These results seem to be consistent with some experiments made at the Lawrence Livermore Laboratory[8] on absorption as a function of laser irradiance. In Ref. 8, a parylene disk was irradiated with a Nd laser pulse of between 5 and 15 Joules with a duration of between 50 to

150 picoseconds. The fractional absorption as a function of irradiance is shown in Fig. 10a of Ref. 8. At the highest irradiance, fractional absorption is greater than fifty percent. Since the irradiance in Fig. 10a of Ref. 8 varies by more than two orders of magnitude, while the irradiance at focus can only vary by at most a factor of nine, much of the variation in irradiance must correspond to a variation in target position with respect to the focal plane. Thus a reasonable conclusion appears to be that in the Livermore experiment,[8] absorption at best focus is around fifty percent, consistent with the other experiments we have cited.

We will now discuss absorption mechanisms currently invoked to explain these data. The favorite mechanism now is probably resonant absorption.[16–22] If the laser light wave vector is not perpendicular to the target, and the electric field is polarized in the plane of incidence (*p* polarization), then the electric field perpendicular to the plane of the target can tunnel from the point of reflection through to the critical density and excite an electrostatic plasma wave. Since this electrostatic wave cannot propagate out, this energy is ultimately deposited into the plasma.

Recent calculations show that the absorption as a function of target tilt angle maximizes at about fifty percent.[21,22] If the electric field is polarized perpendicular to the plane of incidence however (*s* polarization), there is no resonant absorption. It might be possible to explain (with resonant absorption) experiments which show that fractional absorption is not strongly dependent on tilt angle or polarization. One posibility is that laser beams focused by real focusing lenses subtend a substantial range of angles of incidence near focus. One might also argue that the critical surface could be somewhat wrinkled. Experiments,[23] two dimensional simulations[24–26] (for instance see the figures of Ref. 25) and theory[24] all indicate this possibility. Thus the angle and polarization dependence could be partially averaged out. The point to keep in mind however, is that each effect which averages out the angle and polarization dependence will probably also reduce the overall absorption efficiency. For instance, no

matter how the critical surface is wrinkled, at each point on it, the laser light has two possible transverse polarizations; one with the electric vector parallel to the plane of the surface, and one with a component of the electric vector perpendicular to the plane of the surface. Only the latter polarization can give rise to resonant absorption. Thus while a wrinkled critical surface may completely blur the distinction between s and p polarization, the fractional absorption will be cut in half. Furthermore, averaging out the angular dependence will reduce the overall absorption efficiency further still. Thus, if some mechanism (for instance surface wrinkling), averages out the angle and polarization dependence of resonant absorption, the absorption efficiency will probably be reduced from about 50% to about 15%-20%. Clearly this is too low, by at least a factor of two to four, to explain the numerous experiments we have cited.

Recently, experiments at Livermore have conclusively shown that resonant absorption does in fact play an important role in the total absorption picture.[27] With the use of a box calorimeter, very accurate measurements of fractional absorption as a function of angle and polarization could be made. As shown in Fig. 4 of Ref. 27, for s polarization, the fractional absorption at θ (tilt angle) $= 0°$ was about 30%. Fractional absorption was quite independent of θ, it fell gradually to about 15% at $\theta = 60°$. As shown in Fig. 5 of Ref 27, for p polarization, the same qualitative features were observed for roughly $0 < \theta < 20°$ and $45° < \theta < 60°$. However between about 20° and 45° the fractional absorption for p polarization has a maximum. For instance, at $\theta = 30°$, the fractional absorption for p polarization varies between about 35% and 45%, while for s polarization it varies from about 25% to slightly over 30%. Thus resonant absorption is clearly playing an important role.

However even this experiment appears to show that other absorption mechanisms are present. One obvious question is what is causing such 30% absorption for s polarized light. The fact that there is such a distinct maximum at $\theta = 30°$ for p polarized light seems to argue

against large scale averaging out of angles or polarizations. One could in fact argue from this data, that there is some underlying absorption mechanism which is nearly independent of polarization and angle, that is roughly the experimental curve for *s* polarization. However if conditions are right, the absorption can further be enhanced by resonant absorption.

Thus our conclusion, based upon our examination of the generally available data, is that while resonant absorption does play a significant role in the overall picture, there is a clear and pressing theoretical need for some other mechanism also.

The other theoretical mechanisms discussed up to now have been inverse bremsstrahlung and parametric instabilities.[28–31] Both of these mechanisms are faced with severe and possibly catastrophic difficulties in explaining light absorption in a laser produced plasma. These have been discussed at some length in the literature[22] and we will touch only briefly upon them.

Inverse bremsstrahlung is simply collisional damping of the laser light as it propagates through the underdense plasma. For intense laser light, temperatures in the underdense plasma are so high that the collisional damping is too small to give significant absorption in the scale lengths relevant to a laser fusion plasma.[32]

Parametric instabilities, either the decay[29] or oscillating two stream[30] can only occur near the critical density. However it is now well established from theory,[33] numerical simulations,[21,22] and microwave plasma experiments[34] that deposition of laser momentum at the critical surface gives rise to a shock like discontinuity at the critical surface whose width is about half of the free space wavelength of the laser light. This short scale length around the critical surface practically rules out parametric instabilities[35,36,37] as an important absorption mechanism in a laser produced plasma.

A simple derivation of the homogeneous threshold for the decay instability can be given. We show in fact that for those short scale lengths near the critical density, which are appropriate for a laser produced plasma, the decay instability is not expected to be excited. The derivation follows closely that given by Klein et al.[37] More rigorous derivations give the same general results.[36] Say that the density gradient (whose scale length is L) is in the x direction and the laser light electric field is in the y direction. Then, electron plasma waves and ion acoustic waves with wave number **q** in the y direction are excited. However they are only excited in the spatial region where the ion acoustic wave, plasma wave and laser light are all in phase.[35] If K is the wave vector of the laser light, and subscripts i and e denote ion acoustic and electron plasma wave respectively, the length of the region of phase coherence, $\mathscr{L}$, is defined roughly as

$$\int_0^{\mathscr{L}} dx \left(q_{ex} + q_{ix} - K_x \right) \simeq 1. \tag{1}$$

Since the density jump scale length is shorter than a laser wavelength, we neglect the variation of K_x in Eq. (1). The dispersion relation for the ion wave is given by

$$\omega_i = q_{yi} C_s \tag{2}$$

where C_s is the sound speed. Thus there is no variation in q_{xi} with space either. The plasma wave dispersion is

$$\omega^2 = \omega_p^2(x) + 3\left(q_{xe}^2 + q_{ye}^2\right) V_e^2 \tag{3}$$

where V_e is the electron thermal velocity. At x = 0, the plasma wave is assumed to have $q_{xe} = 0$. Solving for q_{xe} as a function of x and integrating Eq. (1), we find

$$\mathscr{L} = \left(\frac{3}{2^{2/3}}\right) L^{1/3} \left(\frac{V_e}{\omega_{pe}(x=0)}\right)^{2/3} \tag{4}$$

In order for the inhomogeneous plasma to be unstable, the waves have to grow significantly while they are in the interaction region, or

$$\gamma^2 \tau_e \tau_i > 1 \tag{5}$$

where γ is the growth rate,[29]

$$\gamma^2 = \frac{\omega_{pe} q_y C_s}{8} \left(\frac{V_{os}}{V_e} \right)^2 \tag{6}$$

where V_{os} is the electron oscillatory velocity amplitude in the field of the laser light. In Eq. (5), τ_i (τ_e) is the time the ion (electron) wave spends in the unstable region. Since the ion wave convects out with the flow velocity V,

$$\tau_i \simeq \mathscr{L}/V. \tag{7}$$

The electron wave refracts out of the unstable region. A simple calculation shows that

$$\tau_e \simeq (L\mathscr{L}/3)^{1/2}/V_e. \tag{8}$$

Combining Eqs. (4-8), we find that the condition for instability is

$$\frac{3}{16} \frac{q_y V_s}{V} L \left(\frac{V_{os}}{V_e} \right)^2 > 1 \tag{9a}$$

Assuming k takes on its maximum value, $q_y \sim \omega_{pe}/2V_e$, $V_s \simeq V$ and $L^{-1} \simeq \omega_p/c$ as shown in Ref. 33, the condition for decay instability on the density jump becomes

$$0.1 \left(\frac{c}{V_e} \right) \left(\frac{V_{os}}{V_e} \right)^2 > 1. \tag{9b}$$

For temperature above a few hundred electron volts, this condition cannot be satisfied for $V_{os} < V_e$, so that the decay instability at least is stabilized by the short gradient scale length.

Thus of the three absorption mechanisms generally invoked, resonant absorption is clearly the most viable. However, equally clearly, experiments and theory show that it cannot account for more than about one third of the measured absorption. The problem then is to find the cause for the remaining absorption. One possible approach is to see if there is any way resonant absorption can be enhanced. For instance, it has recently been shown that in a magnetized plasma, resonant absorption can occur at $\theta = 0$ by coupling to an upper hybrid wave

rather than a plasma wave.[38,39] In this paper, we take another approach and attempt to examine a completely different absorption mechanism. It is enhanced inverse bremsstrahlung resulting from self induced ion density fluctuations.[1] Our conclusion is that this is indeed a viable scheme which shows great promise for explaining the high fractional absorption found in many laser plasma experiments.

III ABSORPTION BY SELF INDUCED ION FLUCTUATIONS IN AN UNMAGNETIZED PLASMA

In this section we review and discuss our earlier work on laser light absorption by ion acoustic fluctuations in an unmagnetized plasma.[1] The idea is that as energy is absorbed, the laser light energy flux is converted into electron thermal energy flux. In order to maintain charge neutrality, a return current of low velocity electrons must flow toward the laser. This return current can drive unstable ion acoustic waves which also propagate toward the laser. These ion acoustic waves have two principal effects.

First of all, they impede electron thermal transport, by anomalously reducing the electron mean free path; and secondly, they give rise to enhanced absorption of the laser light as it encounters the ion density fluctuations.[1] The theory of both of these processes have been given in Refs. 1 and 2 and we will not detail it here.

The remainder of this section will discuss the results of these papers. In Ref. 1, the steady state fluid equations were integrated from $x=0$ backwards toward the laser. Coupled to the equation for the fluid quantities (density, velocity, temperature, and heat flux), were steady state equations for the laser light and also for the ion acoustic fluctuations.

Many of the fluid quantities at $x=0$ could be eliminated by assuming that for small positive x, there is a density jump of the type found by Lee et al.[33] (Recall that the laser is at

$x = -\infty$ and the integration is toward negative x). In terms of the density at $x = 0$, the calculations of Lee et al[33] specify the fluid velocity relative to the ion acoustic velocity and also specify the electron oscillating velocity in the laser field relative to the electron thermal velocity. Thus choosing a lower shelf density and laser irradiance is equivalent to specifying the velocity and temperature at $x = 0$. The remaining quantities to determine were ion temperature and electron energy flux Q. The electron energy flux was chosen so that $Q(-\infty) = 0$, and absorption was studied as a function of ion temperature and density at $x = 0$.

To solve for the laser light absorption, it is necessary to know the spectrum of the ion acoustic turbulence not only in wave number, but also in angle. The reason is that the return current propagates anti-parallel to the Poynting vector of the laser light. Thus the unstable ion acoustic waves propagate mostly in this direction. However only a density fluctuation with a component of its wave vector parallel to the laser light electric field give rise to absorption.[40] Thus one must either do a full two dimensional calculation or else one must assume some angular spectrum. We chose the latter course because earlier experience[41] has shown that a full two dimensional calculation is extraordinarily complex, but it adds very little to the basic physics.

The results of our calculations were in many ways very encouraging. We found high absorption which was not strongly dependent on either polarization or target tilt angle. Also the flux limits calculated, $f \equiv Q/nmV_e^3$ were around 0.1 which were reasonably consistent with experiments on layered targets.[42] Furthermore, this absorption mechanism does not tend to produce a nonthermal tail on the electron distribution function, but rather tends to deposit energy into particles whose velocity is $|\Omega/\mathbf{k}|$ where Ω is the laser light frequency and $\mathbf{k}$ is the wave number of the density fluctuation. This velocity turns out to be about two to four times the electron thermal velocity.

We now turn to some of the difficulties associated with this absorption mechanism. Probably the main problem with the mechanism, as formulated in Ref. 1, is that it requires rather long plasma lengths in the underdense region. For instance for the calculation shown in Fig. 1a of Ref. 1, about 60μ of Nd laser produced plasma was required at half the critical density. In a short pulse experiment, it is not at all clear that the plasma can expand so far. Even in a long pulse experiment, 60μ of plasma at half the critical density will almost certainly give rise to strong Brillouin backscatter.[43,44] Thus, before this mechanism can be really an explanation for significant absorption, the necessary scale lengths of the underdense shelf must be reduced.

There are many ways this can be accomplished. In Ref. 1, the ion acoustic turbulence was started out at a very low level. As is apparent from Fig. 1a of Ref. 1, about half of the necessary length was required simply to build up the turbulence. If the turbulence started out at for instance $e\phi/T_e \sim 10^{-1}$ at $x = 0$, the length required for absorption would be cut in half.

There are many reasons why $e\phi/T_e$ could be this large at $x = 0$. First of all, if there is any transport inhibition on the high density side of the critical surface, the ion acoustic fluctuations responsible for it could propagate to the low density side, initializing the turbulent fluctuations there at some large value. Secondly, the critical surface itself might be unstable. There have been several theoretical and numerical investigations of instabilities which wrinkle the critical surface.[24-26] Also, experiments at NRL[23] show that the critical surface is wrinkled with a mean angle of the wrinkles of about 15°. This wrinkling of the critical surface could certainly initialize ion density fluctions of $e\phi/T_e$ $(=\delta n/n) \sim 0.1$. Thirdly, many authors have speculated that Langmuir solitons could be excited at the critical surface of a laser produced plasma.[45,46] This also could initiate ion density fluctuations there. Finally, it is worth pointing out that the solutions of Lee et al[33] show that the low density shelf is not a uniform plasma; rather it is a plasma with a large amplitude, non sinusoidal density fluctuation

whose wave length is about one half of the free space laser light wavelength. Surely, this large amplitude density fluctuation could also initiate ion acoustic waves at $x = 0$. Thus it is reasonable to assume that the long region in which the turbulence is building up will be eliminated. This reduces the necessary length by about a factor of two, say from 60μ to 30μ.

The length required can be still further reduced by reducing the absorption. For instance if the boundary condition on Q were imposed at $x = -15\mu$ rather than at $x = -\infty$, the problem could be solved the same way, but smaller values of $Q(x = 0)$ (and therefore of fractional absorption) would be found. For instance reducing the length by a factor of two, from 30μ to 15μ would probably just reduce the fractional absorption by a factor of roughly two. Since Ref. 1 found very high absorption, typically forty to sixty percent, the fractional absorption can be reduced by a factor of two and this will still be a very important absorption mechanism. Thus we conclude that length scales can be greatly reduced from those shown in Ref. 1 while absorption remains high.

Another possible problem is that the mechanism in Ref. 1 only works, for the boundary conditions imposed, if T_e/T_i is larger than about 15. The reason is that as the ion temperature increases, enough ions have velocity $\simeq \alpha\ (T_i/M)^{1/2}$ where $\alpha \geqslant 1$, comparable to the wave phase velocity $(T_e/M)^{1/2}$ so that ion Landau damping is important and the unstable region is too small. While temperatures ratios as high as 15 to 20 are not unreasonable, and have in fact been observed in fluid simulations of laser produced plasmas,[47] the absorption mechanism would be more universal if this temperature ratio could be reduced. We find that the temperature ratios for efficient absorption can in fact be significantly reduced if the ion charge $Z > 1$. The reason is that the wave phase velocity increases to $(ZT_e/M)^{1/2}$ while the ion velocity remains $\alpha\ (T_i/M)^{1/2}$ so that the importance of ion Landau damping is reduced. In a later section, we show that if $Z = 2.7$, corresponding roughly to a CH_2 plasma, then there can be absorption with temperature ratios as low as five.

A final problem with this absorption mechanism is that it depends on a number of parameters which cannot be a-priori specified. These include the lower shelf density, the ion temperature, and the boundary condition on Q at $x = -\infty$. In a magnetized plasma, the situation is even worse, because the magnetic field must also be specified. In a later section we will show how these parameters and boundary conditions can be specified by solving for the dynamics of the problem rather than the steady state.

IV ABSORPTION IN A MAGNETIZED PLASMA, QUALITATIVE DISCUSSION

It is now well established by both theory[48-54] and experiment[55] that magnetic fields of order a megagauss can be generated in a laser produced plasma. Thus there is excellent motivation for studying anomalous absorption by ion acoustic waves in a magnetized plasma.

In a magnetized plasma the absorption process discussed in the previous section can be enhanced by a very considerable amount as will be discussed shortly. Let us for now assume slab geometry so that the magnetic field is in the z direction, and gradients are in the x direction. Then, it was shown in Ref. 3 that ion acoustic waves which propagate principally in the y direction are driven unstable. The growth rate is proportional to dT_e/dx and the electron thermal energy flux in the x direction is proportional to $(e\phi/T_e)^2 \, dT_e/dx$.

Fluid equations, equations for the ion acoustic turbulence, and equations for the laser light are integrated back from $x = 0$ to $x = -\infty$. For small positive x, a density jump of the type found by Lee et al[33] is once again assumed. Thus the fluid velocity and temperature can be determined in terms of lower shelf density and laser irradiance. We have found that a different boundary condition on $Q(x = -\infty)$ must be imposed. The reason is that if Q starts out positive at $x = 0$, it always ends up negative at $x = -\infty$ no matter how parameters are varied. That is, Q at $x = -\infty$ can never be zero.

Actually an outward energy flux for large negative x may even be a more reasonable boundary condition than $Q(x = -\infty) = 0$ as assumed in Ref. 1. For instance if the plasma for negative x expands as a rarefaction wave,

$$n = n_o \exp\left[\frac{x}{C_s t}\right] \tag{10a}$$

$$V = -C_s - \frac{x}{t} \tag{10b}$$

where $C_s = (Z T_e/M)^{1/2}$, then the fluid energy flux to the left at $x = 0$ is

$$W_f = -1/2 n_o M C_s^3 - 5/2\, n_o Z C_s T_e = -3 n_o M C_s^3 . \tag{11}$$

However one can integrate Eqs. (10) to get the rate of change of energy in the region $x < 0$. The energy in this region increases faster than the energy flux (Eq. (11)) times the area. Thus additional outward energy flux is needed to feed the isothermal rarefaction wave. The total incident energy flux needed is $-4 n_o M C_s^3$. Thus if the outward electron thermal energy flux at $x = -\infty$ is $-n_o M C_s^3$, the total outward energy flux will be just sufficient to propagate an isothermal rarefaction wave. Thus the boundary condition we assume is

$$Q(x = -\infty) = -n_o M C_s^3 \tag{12}$$

Another difference between the calculations presented in the next section and those in Ref. 1 is that we generally assume 20% resonant absorption at the density jump. (However a few calculations were done without resonant absorption as a comparison). One advantage of this absorption mechanism is that it does not compete with resonant absorption but co-exists very easily with it. This contrasts with the case for parametric instabilities. Parametric instabilities and resonant absorption can occur only at the critical density. However steepening at the critical density tends to enhance resonant absorption and reduce the effect of parametric instabilities. Broadening the profile at the critical density does the opposite. Thus resonant absorption and parametric instabilities compete with each other.

Absorption by self generation of ion turbulence and resonant absorption do not compete with each other however, since they occur at different locations in the plasma. If the density profile is steep near the critical surface, but broad in the underdense region as shown in Fig. 1, then both processes can occur simultaneously.

There are several reasons why this absorption mechanism can be enhanced in a magnetized plasma. First of all for a given $Q(x=0)$; the smaller $e\phi/T_e$ is, the larger $\frac{dT_e}{dx}$ is, and therefore the larger the growth rate is. Thus, even if $e\phi/T_e$ is assumed very small at $x=0$, the growth rate is very large, and the transient region of small $e\phi/T_e$ is effectively eliminated.

Secondly, if the laser light is polarized with its electric field vector **E** in the y direction, it and the wave numbers, **k**, of the ion acoustic turbulence are aligned for maximum absorption. Thus there is no need for making any assumptions at all about the angular spectrum of the turbulence. In fact a simple calculation shows that the same $e\phi/T_e$ and $|\mathbf{k}|$ gives rise to about 6 times as much absorption if **k** and **E** are aligned than it does for the angular spectrum assumed in Ref. 1. Thus here we find roughly equivalent absorption, but with much smaller values of turbulent field strength and scale length. (Of course if **E** is in the z direction, then one must again make assumptions about the angular width of the spectrum.)

Thirdly, the ion acoustic wave propagates in the y direction rather than the negative x direction. Thus its group velocity in the x direction is just the fluid speed, rather than the fluid speed plus the sound speed. Therefore the spatial growth rate γ/V_{gr} is larger in the magnetized plasma case even if γ is unchanged. To summarize, the effect of a magnetic field can significantly improve the absorption of laser light by self generation of ion acoustic turbulence.

We conclude this section with some speculations on the geometry of the magnetic field at the critical surface. Let us take the case of a focused laser beam striking a slab target. Then the geometry is basically cylindrical; the magnetic field lines are circles as shown in Fig. 2a.

The polarization of the laser light can be either up and down (s polarization) or right to left (p ploarization). It is not clear whether the temperature gradient is principally axial or principally radial. If $dT_e/dz > dT_e/dr$, one can draw some interesting conclusions. The regions above and below the center of the laser spot have $\mathbf{E} \perp \mathbf{B} \perp \mathbf{k}$, leading to maximum absorption there, if the light has s polarization. (For p polarization, maximum absorption is to the right and left of the laser spot.)

Now consider what happens if the target is tilted. The geometry no longer has cylindrical symmetry and the magnetic field lines will be distorted into ellipses, as shown in Fig. 2b. Notice now that $\mathbf{E} \perp \mathbf{B} \perp \mathbf{k}$, on the long side of the ellipse for s polarization, while $\mathbf{E} \perp \mathbf{B} \perp \mathbf{k}$ on the short side of the ellipse for p polarization. Thus absorption by self generated ion acoustic turbulence favors s polarization over p polarization if the target is tilted. This may be one possible explanation for the fact that even in experiments which show resonant absorption,[27] the absorption of p polarized light is only about 50% higher than the corresponding absorption for s polarized light.

V ABSORPTION IN A MAGNETIZED PLASMA

QUANTITATIVE CALCULATIONS

In this section we describe in more detail the calculations of absorption efficiency in a magnetized plasma. The steady state fluid is described by the four equations for number density, total momentum, electron energy and ion energy. They are

$$nV = \text{constant} \tag{13a}$$

$$\frac{n}{Z} MV \frac{dV}{dx} = -\frac{d}{dx} n\,(T_e + T_i/Z) - \frac{\omega_{pe}^2}{2\Omega^2} \frac{d}{dx} \left(\frac{E_i^2 + E_r^2}{8\pi} \right) \tag{13b}$$

$$\frac{3}{2} nV \frac{dT_e}{dx} + nT_e \frac{dV}{dx} = -\frac{d}{dx} Q_e + C_{Te} + \nu_{an} \left(\frac{\omega_{pe}}{\Omega^2} \right) \left(\frac{E_i^2 + E_r^2}{8\pi} \right) \tag{13c}$$

$$\frac{3}{2}\frac{n}{Z}V\frac{dT_i}{dx} + \frac{n}{Z}T_i\frac{dV}{dx} = -C_{Te}. \tag{13d}$$

Above, n is the electron number density, E_i^2 (E_r^2) is the electric field squared of the incident (reflected) laser light, C_{Te} is the ion acoustic wave induced energy exchanged between electrons and ions, ν_{an} is the ion acoustic wave induced collisional damping and Q_e is the electron thermal flux. These three quantities will be specified shortly.

The equations for the spatial evaluation of the ion acoustic waves are

$$\frac{d}{dx}\left|\frac{e\phi(k)}{T_e}\right|^2 = \frac{2\gamma(k)}{V}\left|\frac{e\phi(k)}{T_e}\right|^2 \tag{14}$$

where ϕ is the electrostatic potential at wave number k, V is the group velocity, which in this case is the fluid velocity, and $\phi(k)$ is the growth rate[3] of a mode at wave number $\mathbf{k}$

$$\gamma(k) = \left(\frac{\pi m}{8M}\right)\left[\frac{3}{2}\frac{kc}{eB}\frac{dT_e}{dx} - |k|\left(\frac{ZT_e/M}{1+k^2\lambda_D^2} + \frac{3T_i}{M}\right)^{1/2}\right.$$

$$\left.\times\left\{1 + Z\left(\frac{M}{m}\right)^{1/2}\left(\frac{T_e}{T_i}\right)^{3/2}\exp -\left(\frac{ZT_e}{2T_i(1+k^2\lambda_D^2)} + \frac{3}{2}\right)\right\}\right] \tag{15}$$

where B is the magnetic field. Equation (15) above is the growth rate given in Ref. 3, but with ion Landau damping added on and also with account taken of the fact that $Z \neq 1$. The mode phase velocity is assumed to be in the y direction and the gradients are assumed to be in the x direction. The configuration we adopt has the laser at $x = -\infty$ and the critical surface just to the right of $x = 0$. Since energy flux is to the right, the plasma temperature gradient is negative. Therefore, as is apparent from Eq. (15), unstable waves propagate in the negative y direction.

Coupled to the equations for fluid quantities and ion acoustic waves are the equations for the incident and reflected laser light. If θ is the target tilt angle, these equations are

$$c\frac{d}{dx}\left(\cos^2\theta - \frac{\omega_p^2(x)}{\Omega^2}\right)E_i^2 = -\nu_{an}\frac{\omega_{pe}^2}{\Omega^2}E_i^2 \tag{16a}$$

$$c\frac{d}{dx}\left(\cos^2\theta - \frac{\omega_p^2(x)}{\Omega^2}\right)E_r^2 = \nu_{an}\frac{\omega_{pe}^2}{\Omega^2}E_r^2 \tag{16b}$$

The remaining quantities to specify are Q, C_{Te} and ν_{an}. The quantity Q is given in Ref. 3 and is

$$Q = -\frac{13}{4}\left(\frac{\pi}{2}\right)^{1/2}\sum_k\left|\frac{e\phi(k)}{T_e}\right|^2\frac{k^2}{|k|}\frac{V_e^3}{\omega_{ce}^2}n\frac{dT_e}{dx}, \tag{17}$$

where ω_{ce} is the electron cyclotron frequency. The quantity C_{Te} is given in Ref. 2 for an unmagnetized plasma. Using Ref. 3, C_{Te} can easily be generalized to the case of a magnetized plasma. The result is

$$C_{Te} = \left(\frac{\pi}{8}\right)^{1/2}\sum_k\left|kC_s\right|\left(nmC_sV_e - \frac{3nV_e\,mck}{2eB|k|}\frac{dT_e}{dx}\right)\left|\frac{e\phi(k)}{T_e}\right|^2. \tag{18}$$

Notice that the second term in the parentheses above cools electrons since $k, \frac{dT_e}{dx} < 0$. The explanation for this effect is given in Ref. 2. The one remaining thing to specify is ν_{an}. This is calculated as in Refs. 1 and 2, using the theory of Dawson and Oberman.[40] However now the absorption is enhanced because in the configuration we examine, $\mathbf{k} \parallel \mathbf{E}$. The result is

$$\nu_{an} = \frac{\omega_{pe}^2}{\Omega}\sum_k\frac{2\left(\frac{\omega_p}{\Omega}\right)^3|k|\lambda_D}{\left[1-\left(\frac{\omega_p}{\Omega}\right)^2\right]^2 + 4\left(\frac{\omega_p}{\Omega}\right)^6k^2\lambda_D^2}\left|\frac{e\phi(k)}{T_e}\right|^2 \tag{19}$$

where λ_D is the electron Debye length. The fluid system is now completely specified.

Equations (13)-(19) are integrated numerically back towards $x = -\infty$, subject to the boundary conditions described in the previous section. Also $\sum_k\left|\frac{e\phi(k)}{T_e}\right|^2 = 10^{-3}$ at $x = 0$. Typical solutions are shown in Figs. 3 a and b for a Nd laser produced plasma. In both cases

$T_e/T_i = 20$, $B = 10^6$ Gauss, $n(x=0) = 1/2\ n_{cr}$, $\theta = 0$, and the laser irradiance is $2 \times 10^{15}\ W/cm^2$. In Fig. 3a no resonant absorption is assumed to the right of $x=0$ and the total fractional absorption is 16%. In Fig. 3b, 20% resonant absorption is assumed, and the fractional absorption is 33%. In each case, 10 values of k were used. Notice that in each case the absorption takes place in a length of less than 10 microns. Also the turbulent field strength, $\left|\frac{e\phi}{T_e}\right|^2 < 10^{-2}$, is less than that calculated in Ref. 1.

In Fig. 4 is shown a plot of flux limit $f = Q(x=0)/nmV_e^3 \cdot T_e(x=0)$ and fractional absorption as a function of laser irradiance between about $4 \times 10^{14}\ W/cm^2$ and $1.5 \times 10^{16}\ W/cm^2$. Again $n(x=0) = 1/2\ n_{cr}$, $T_e/T_i = 20$, $\theta = 0$, $B = (1\ \text{MG}) \times (T_e(x=0)/12\ \text{keV})^{1/2}$. This functional form for B specifies a constant plasma β as the temperature is increased. Twenty percent resonant absorption at the critical surface is assumed here. Actually the flux limit f should be somewhat higher than shown here, because the electron thermal flux arising from the 20% assumed resonant absorption was not included in $Q(x=0)$. Figure 4 indicates the same general result found in Ref. 1, the fractional absorption and flux limits are weakly varying functions of laser irradiance, while the electron temperature at $x=0$ is a strongly varying function of it.

In another series of runs, we tested the dependence of fractional absorption on magnetic field strength. For $n(x=0) = 1/2\ n_{cr}$, $\theta = 0$, $T_e/T_i = 20$, and assumed resonant absorption of 20% and a laser irradiance of $2 \times 10^{15}\ W/cm^2$, the magnetic field was varied from 700 KG to 2.2 MG. As is shown in Fig. 5, the relative absorption is about 40% nearly independent of field strength. However the required length for absorption increases as the field strength decreases. Thus if length scale is an important feature in limiting the absorption, higher magnetic fields give rise to higher absorption.

In another series of calculations, the density at $x=0$ was varied assuming, $\theta=0$, $I = 2 \times 10^{15}\ W/cm^2$, $T_e/T_i = 20$, $B = 1\ MG$, and 20% resonant absorption at $x=0$. The

result is shown in Fig. 6. The results are quite similar to earlier calculations in a field free plasma. At $n(x=0)=0.25\ n_{cr}$, there is virtually no absorption (recall 20% resonant absorption was assumed). The fractional absorption however is a steeply rising function of density at $x=0$. Notice that the fractional absorption rises to almost 80% for $n(x=0)=0.7n_{cr}$.

In a final series of calculations we examined the functional dependence of fractional absorption on temperature ratio for $n(x=0)=0.5n_{cr}$, $I=2.10^{15}\ W/cm^2$, $B=1MG$ and also assuming 20% resonant absorption. The results are shown in Fig. 7. As is apparent, one finds absorption at temperature ratios as low as 5. One very interesting and surprising result is the maximum for low temperature ratios. We find a qualitatively different behavior at low temperature ratio from that at high temperature ratio. At high temperature ratio, the turbulent spectrum is dominated by a single k at all x. For low temperature ratios however different k's dominate the spectrum at different positions. This gives rise to a choppy behavior of $n(x)$ and $\frac{e\phi}{T_e}(x)$ as shown in Fig. 8.

VI SCALING WITH LASER WAVELENGTH

All of the calculations in the previous section were done assuming the laser wavelength was 1.06μ. However it is a simple matter to show that these equations have a simple scaling so that a solution for one laser wavelength can be scaled into a solution for any laser wavelength. Let us assume that the laser wavelength increases by a factor of α, then the equations in the previous section are all invariant under the following scale transformation

$$x \rightarrow \alpha x$$

$$n \rightarrow n/\alpha^2$$

$$T \rightarrow T$$

$$V \rightarrow V$$

$$E^2 \rightarrow E^2/\alpha^2$$

$$\Omega \rightarrow \Omega/\alpha$$

$$B \rightarrow B/\alpha$$

$$\frac{e\phi}{T_e} \rightarrow \frac{e\phi}{T_e}. \tag{20}$$

Some of the results of this scaling transformation are

$$Q \rightarrow Q/\alpha^2$$

$$\nu_{an} \rightarrow \nu_{an}/\alpha$$

$$C_{Te} \rightarrow C_{Te}/\alpha^3$$

$$\gamma \rightarrow \gamma/\alpha \tag{21}$$

By examining the equations for the unmagnetized plasma in Ref. 1, it is not difficult to see that these equations are invariant under the same scale transformation. Thus Figs. 3-8 can easily be interpreted for a CO_2 laser produced plasma as well as for a Nd laser produced plasma.

VII THE DYNAMICS OF THE INTERACTION

The calculations we have performed in this paper and in our previous work[1] are all steady state, that is all $\partial/\partial t$'s are set equal to zero in all fluid equations. As we have discussed previously, one drawback with this approach is that several quantities have to be specified arbitrarily. One possible way to specify these quantities in a physical way is to solve for the dynamics of the interaction[56] rather than the steady state. The way to include anomalous transport and absorption into a dynamic one or two dimensional fluid simulation is given in Ref. 2 for an unmagnetized plasma, and in Ref. 3 for a magnetized plasma. By solving the dynamic problem, one could then solve for such quantities as $Q(x = -\infty)$, $n(x = 0)/n_{cr}$, etc. in terms of boundary conditions imposed far from the critical surface, that is in the solid target and in the vacuum.

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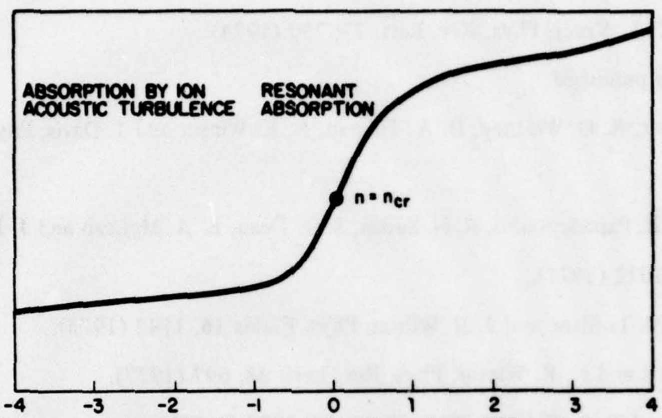

Fig. 1 — Schematic diagram indicating where different absorption mechanisms occur. Unit on the horizontal axis is free space wavelength of laser light.

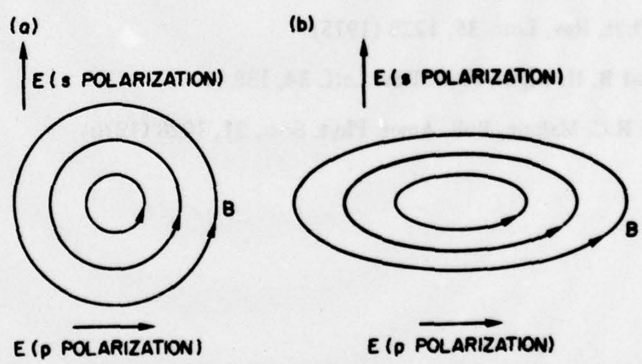

Fig. 2 — Magnetic field lines and various polarizations for the case of a) target normal to the incident laser light, b) tilted target.

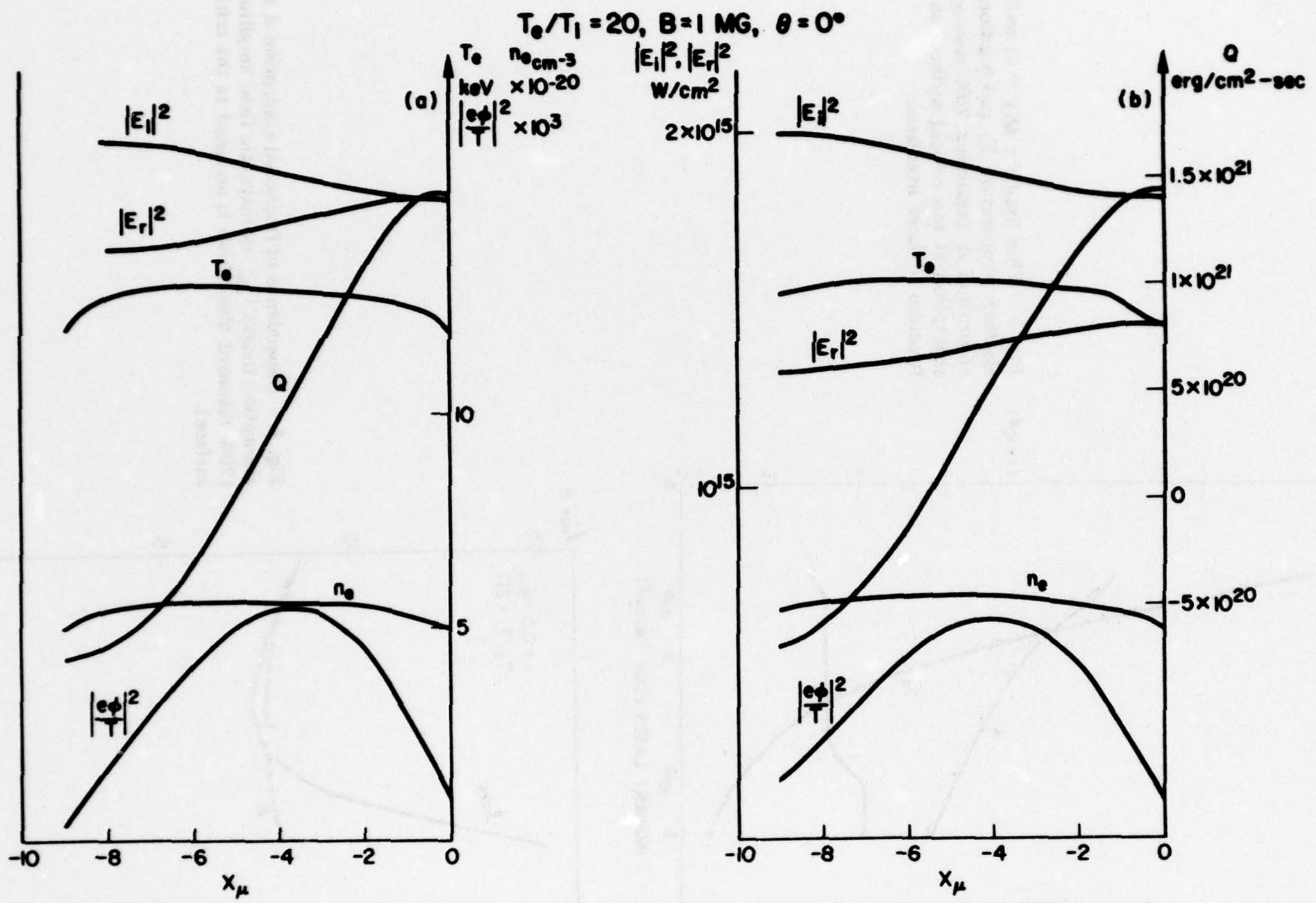

Fig. 3 — Typical solution for the case of a) no resonant absorption beyond $x = 0$, b) 20% resonant absorption beyond $x = 0$. Note that $x = 0$ corresponds, in this case, to $n = 0.5n_{crit}$.

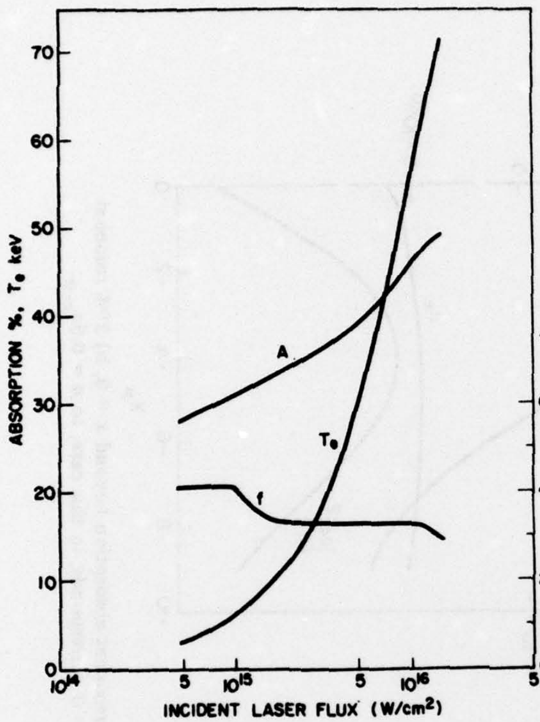

Fig. 4 — Flux limit $f = Q(x = 0)\, nmV_e^3$, electron temperature T_e and fractional absorption A (assuming 20% resonant absorption at the critical surface) as a function of laser irradiance.

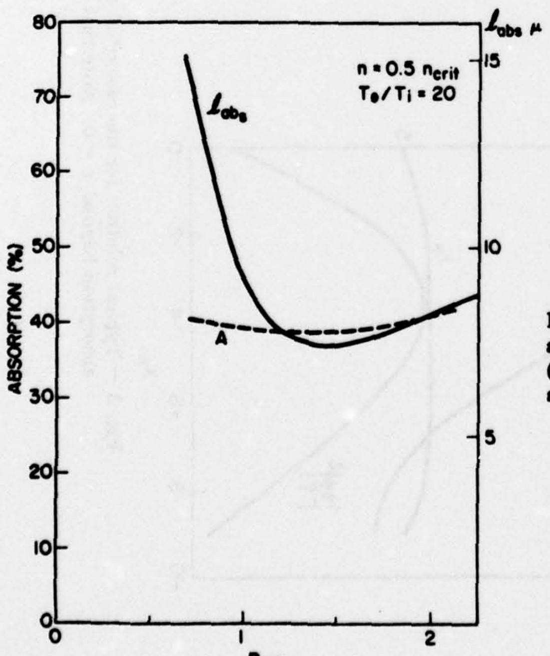

Fig. 5 — Dependence of fractional absorption A and absorption length l_{abs} on magnetic field amplitude. (20% resonant absorption is assumed at the critical surface).

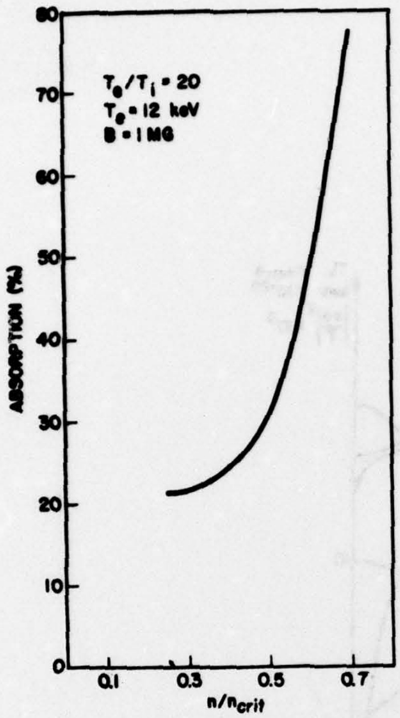

Fig. 6 — Fractional absorption as a function of the lower shelf density n/n_{crit}.

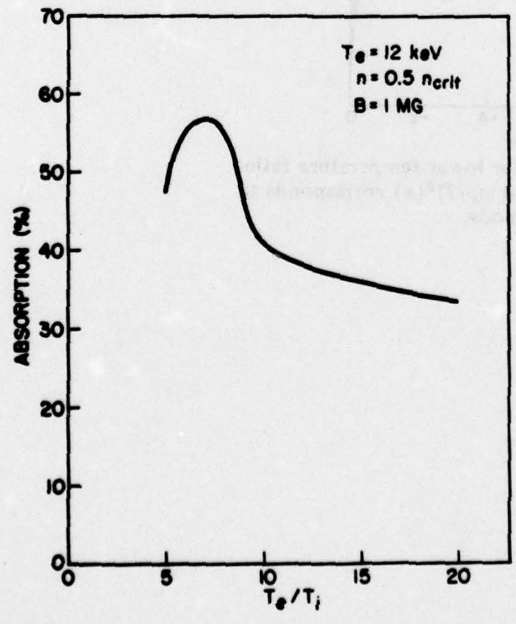

Fig. 7 — Fractional absorption as a function of electron to ion temperature ratio for the case of Fig. 3b).

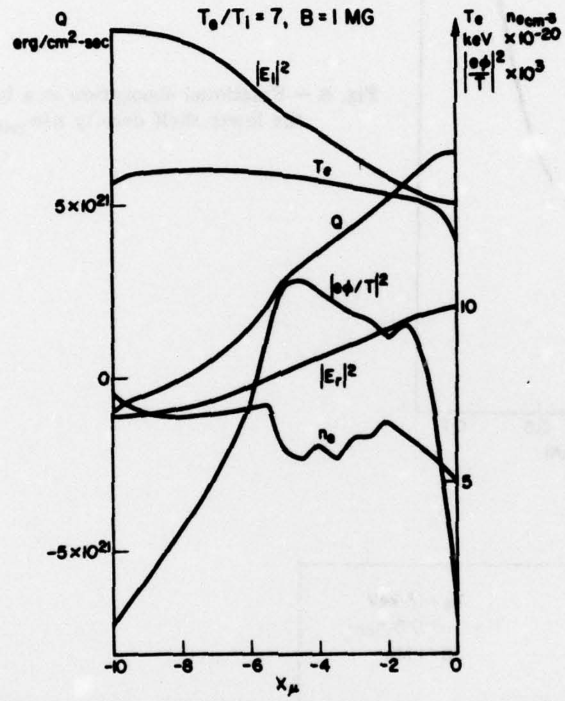

Fig. 8 — Typical solution for lower temperature ratio. Choppy behavior of $n_e(x)$ and $|e\varphi/T|^2(x)$ corresponds to shifting dominant unstable mode.